

The Standard-smooth Variant of Hybrid Inflation

George Lazarides

Physics Division, School of Technology, Aristotle University of Thessaloniki, Thessaloniki 54124, Greece

Abstract.

We consider the extension of the supersymmetric Pati-Salam model introduced in order to solve the b -quark mass problem in supersymmetric theories with Yukawa unification, universal boundary conditions and $\mu > 0$. This model naturally leads to the new shifted and new smooth hybrid inflation scenarios, which, however, yield, in minimal supergravity, too large values of the spectral index n_s . We show that this problem can also be resolved within the same model by a two-stage inflationary scenario based only on renormalizable superpotential interactions. The first stage is of the standard and the second of the new smooth hybrid type. The cosmological scales exit the horizon during the first stage of inflation and acceptable n_s 's can be achieved by restricting the number of e-foldings of our present horizon during this inflationary stage. The additional e-foldings needed for solving the horizon and flatness problems are naturally provided by the second stage of inflation. Monopoles are formed at the end of the first stage of inflation and are, subsequently, diluted by the second stage of inflation so that their density in the present universe is utterly negligible.

Keywords: Inflation, Supersymmetry

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1. INTRODUCTION

One of the most promising inflationary scenarios is hybrid inflation (HI) [1], which is naturally realized [2, 3] in supersymmetric (SUSY) grand unified theory (GUT) models. In its standard realization, though, the GUT gauge symmetry (G) breaking occurs at the end of inflation leading [4] to a disastrous production of magnetic monopoles if these defects are predicted by the underlying symmetry breaking. This problem is avoided in the smooth [4, 5] or shifted [6] variants of SUSY HI, where G is broken to the standard model gauge group (G_{SM}) already during inflation (for a review, see Ref. [7]).

These two variants of SUSY HI, originally based on non-renormalizable superpotential terms, can be implemented [8, 9] with only renormalizable terms in an extended SUSY GUT model based on the Pati-Salam (PS) gauge group $G_{\text{PS}} = \text{SU}(4)_c \times \text{SU}(2)_L \times \text{SU}(2)_R$ [10], whose breaking to G_{SM} predicts [11] the existence of doubly charged monopoles.

It is very interesting to point out that this extended SUSY PS model was motivated [12] (see also Ref. [13]) by a very different issue: In SUSY models with Yukawa unification (YU)[14], such as the simplest SUSY PS model (see Ref. [15]), and universal boundary conditions, the mass of the b -quark (m_b) turns out [16] to be too large for $\mu > 0$. By appropriately extending the model, however, YU is modestly violated and m_b can be adequately reduced.

Fitting the Wilkinson microwave anisotropy probe (WMAP) data [17] with the standard power-law cosmological model with cold dark matter and a cosmological constant, one obtains a spectral index n_s clearly lower than unity. However, in canonical supergravity (SUGRA), these HI models yield [18] n_s 's which are very close to unity or even larger than unity.

One way [19] to reduce n_s is by restricting the number of e-foldings of our present horizon during the main part of inflation which is responsible for the observed density perturbations. The additional e-foldings required for solving the horizon and flatness problems can be provided by a subsequent second stage of inflation.

We show [20] that the same extended SUSY PS model of Ref. [12] leads to a two-stage inflationary scenario with acceptable values of n_s in minimal SUGRA. The first stage is of the standard hybrid type, while the second is of the smooth hybrid type. So, the name standard-smooth HI is coined for this scenario. Standard HI occurs along a trivial flat direction on which G_{PS} is unbroken. As the inflaton drops below a critical value, this direction is destabilized giving its place to a non-flat valley for smooth HI on which G_{PS} is broken.

Note that the same extended SUSY PS model can lead [21] to an alternative inflationary scenario with cosmic strings [22] (for a textbook presentation or a review, see e.g. Ref. [23]), since it possesses a semi-shifted flat direction on which $U(1)_{B-L}$ is unbroken. When the system crosses the critical point of this path, $U(1)_{B-L}$ breaks and cosmic strings are produced contributing to the primordial perturbations, in which case, the data allow [24] larger values of n_s . In this semi-shifted inflationary scenario, no magnetic monopoles are formed.

2. THE EXTENDED SUSY PS MODEL

In this model [12] (see also Refs. [6, 13]), the breaking of G_{PS} to G_{SM} is achieved by the vacuum expectation values (VEVs) of a conjugate pair of Higgs superfields H^c and \bar{H}^c belonging to the $(\bar{4}, 1, 2)$ and $(4, 1, 2)$ representations of G_{PS} respectively. There also exist a gauge singlet S and a conjugate pair of superfields $\phi, \bar{\phi} \in (15, 1, 3)$ with $\langle \phi \rangle$, the VEV of ϕ , breaking G_{PS} to $G_{SM} \times U(1)_{B-L}$. In addition to G_{PS} , the model possesses a Z_2 matter parity symmetry and two global $U(1)$ symmetries, namely a Peccei-Quinn and a R symmetry.

The superpotential terms relevant for inflation are [9]

$$W = \kappa S(M^2 - \phi^2) - \gamma S H^c \bar{H}^c + m \phi \bar{\phi} - \lambda \bar{\phi} H^c \bar{H}^c, \quad (1)$$

where the mass parameters M, m ($\sim M_{GUT}$, the SUSY GUT scale) and any two of the three dimensionless coupling constants κ, γ, λ can be made real and positive by an appropriate rephasing of the fields. We choose the third dimensionless coupling constant to be real and positive too.

The F-term scalar potential obtained from the superpotential W in Eq. (1) is

$$V = |\kappa(M^2 - \phi^2) - \gamma H^c \bar{H}^c|^2 + |m\bar{\phi} - 2\kappa S\phi|^2 + |m\phi - \lambda H^c \bar{H}^c|^2 + |\gamma S + \lambda \bar{\phi}|^2 (|H^c|^2 + |\bar{H}^c|^2). \quad (2)$$

From this potential and the vanishing of the D-terms (which yields $\bar{H}^{c*} = e^{i\theta} H^c$), we find [9] two distinct continua of SUSY vacua:

$$\phi = \phi_+, \quad \bar{H}^{c*} = H^c, \quad |H^c| = \sqrt{m\phi_+/\lambda} \quad (\theta = 0); \quad (3)$$

$$\phi = \phi_-, \quad \bar{H}^{c*} = -H^c, \quad |H^c| = \sqrt{-m\phi_-/\lambda} \quad (\theta = \pi) \quad (4)$$

with $\bar{\phi} = S = 0$, where

$$\phi_{\pm} \equiv \frac{\gamma m}{2\kappa\lambda} \left(-1 \pm \sqrt{1 + \frac{4\kappa^2\lambda^2 M^2}{\gamma^2 m^2}} \right). \quad (5)$$

3. FLAT DIRECTIONS

The potential V in Eq. (2) possesses [9] three flat directions:

- (i) The trivial flat direction: $\phi = \bar{\phi} = H^c = \bar{H}^c = 0$ with $V = V_{\text{tr}}^0 \equiv \kappa^2 M^4$.
- (ii) The new shifted one (for $\gamma \neq 0$) on which G_{PS} is broken to G_{SM} :

$$\begin{aligned} \phi &= -\frac{\gamma m}{2\kappa\lambda}, \quad \bar{\phi} = -\frac{\gamma}{\lambda} S, \quad H^c \bar{H}^c = \frac{\kappa\gamma(M^2 - \phi^2) + \lambda m\phi}{\gamma^2 + \lambda^2}, \\ V &= V_{\text{nsh}}^0 \equiv \frac{\kappa^2\lambda^2}{\gamma^2 + \lambda^2} \left(M^2 + \frac{\gamma^2 m^2}{4\kappa^2\lambda^2} \right)^2. \end{aligned} \quad (6)$$

- (iii) The semi-shifted one (for $\tilde{M}^2 \equiv M^2 - m^2/2\kappa^2 > 0$):

$$\phi = \pm \tilde{M}, \quad \bar{\phi} = \frac{2\kappa\phi}{m} S, \quad H^c = \bar{H}^c = 0 \quad (7)$$

on which G_{PS} is broken to $G_{\text{SM}} \times \text{U}(1)_{\text{B-L}}$ ($\phi \neq 0$, $H^c = \bar{H}^c = 0$) and $V = V_{\text{ssh}}^0 \equiv \kappa^2(M^4 - \tilde{M}^4)$.

4. STANDARD-SMOOTH HI

We will take $\tilde{M}^2 < 0$. In this case, the semi-shifted flat direction does not exist. Also, $V_{\text{nsh}}^0 > V_{\text{tr}}^0$ and, thus, the system will eventually settle down on the trivial flat direction. Expanding ϕ , $\bar{\phi}$, H^c , \bar{H}^c as $s = (s_1 + i s_2)/\sqrt{2}$, we find [20], on the trivial direction, the mass² matrices $M_{\phi_1}^2$ of ϕ_1 , $\bar{\phi}_1$ and $M_{\phi_2}^2$ of ϕ_2 , $\bar{\phi}_2$:

$$M_{\phi_1(\phi_2)}^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 \mp 2\kappa^2 M^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (8)$$

and the mass² matrices $M_{H_1}^2$ of H_1^c , \bar{H}_1^c and $M_{H_2}^2$ of H_2^c , \bar{H}_2^c :

$$M_{H_1(H_2)}^2 = \begin{pmatrix} \gamma^2|S|^2 & \mp \gamma\kappa M^2 \\ \mp \gamma\kappa M^2 & \gamma^2|S|^2 \end{pmatrix}. \quad (9)$$

The matrices $M_{H1(H2)}^2$ acquire one negative eigenvalue for $|S| < S_c \equiv \sqrt{\kappa/\gamma} M$ and the trivial direction becomes unstable.

For $\gamma \ll 1$, the trivial direction, after its destabilization at S_c , gives its place [9] to a valley of minima with $\theta \simeq 0$, which leads to the corresponding SUSY vacua in Eq. (3). This valley possesses a classical inclination and can accommodate a second stage of inflation of the new smooth hybrid type.

The standard-smooth HI scenario goes [20] as follows:

(i) The system initially inflates on the trivial direction which acquires a slope from the one-loop radiative correction (RC) [3] due to the SUSY breaking caused by the constant potential energy density $V = V_{\text{tr}}^0$. So, we get a first stage of inflation of the standard hybrid type.

(ii) As the system moves below S_c , G_{PS} breaks to G_{SM} . After a short intermediate inflationary phase, the system settles down on the new smooth path and new smooth HI occurs. The second stage of inflation (intermediate phase plus new smooth HI) yields the additional e-foldings for solving the horizon and flatness problems. At the end, the system falls into the SUSY vacua leading, though, to no magnetic monopoles, since G_{PS} is broken to G_{SM} during the second stage of inflation.

However, two important requirements must be fulfilled [20]:

(a) The number of e-foldings during the second stage of inflation must be adequately large for diluting any monopoles generated at the end of the first stage.

(b) Cosmological scales get perturbations only from the first stage of inflation.

4.1. One-loop RC

The one-loop RC [3] to the potential V from SUSY breaking on the trivial path is calculated by the Coleman-Weinberg formula [25]:

$$\Delta V = \frac{1}{64\pi^2} \sum_i (-1)^{F_i} M_i^4 \ln \frac{M_i^2}{\Lambda^2}, \quad (10)$$

where we sum over helicity states, F_i and M_i^2 are the fermion number and mass² of the i th state and Λ is a renormalization scale.

So, we need the mass spectrum of the model on the trivial path. We find [20] two groups of 45 pairs of real scalars with mass² matrices

$$M_{-(+)}^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 \mp 2\kappa^2 M^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (11)$$

and two more groups of 8 pairs of real scalars with mass² matrices

$$M_{1(2)}^2 = \begin{pmatrix} \gamma^2|S|^2 & \mp \gamma \kappa M^2 \\ \mp \gamma \kappa M^2 & \gamma^2|S|^2 \end{pmatrix}. \quad (12)$$

Also, 45 pairs of Weyl fermions with mass² matrix

$$M_0^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (13)$$

and 8 more pairs of Weyl fermions with mass² matrix

$$\bar{M}_0^2 = \begin{pmatrix} \gamma^2 |S|^2 & 0 \\ 0 & \gamma^2 |S|^2 \end{pmatrix}. \quad (14)$$

The one-loop RC to the inflationary potential V is then

$$\begin{aligned} \Delta V = & \frac{45}{64\pi^2} \text{tr} \left[M_+^4 \ln \frac{M_+^2}{\Lambda^2} + M_-^4 \ln \frac{M_-^2}{\Lambda^2} - 2M_0^4 \ln \frac{M_0^2}{\Lambda^2} \right] \\ & + \frac{8}{64\pi^2} \text{tr} \left[M_1^4 \ln \frac{M_1^2}{\Lambda^2} + M_2^4 \ln \frac{M_2^2}{\Lambda^2} - 2\bar{M}_0^4 \ln \frac{\bar{M}_0^2}{\Lambda^2} \right]. \end{aligned} \quad (15)$$

The total effective potential on the trivial path in global SUSY is

$$V_{\text{tr}} = v_0^4 + \Delta V, \quad (16)$$

where $v_0 \equiv \sqrt{\kappa}M$ is the inflationary scale. Note that the $\sum_i (-1)^{F_i} M_i^4$ is S -independent implying that the slope of the path is Λ -independent. This guarantees that the observables do not depend on Λ , which remains undetermined.

4.2. SUGRA correction

The F-term scalar potential in SUGRA is

$$V = e^{K/m_{\text{P}}^2} \left[(F_i)^* K^{i^*j} F_j - 3 \frac{|W|^2}{m_{\text{P}}^2} \right], \quad (17)$$

where K is the Kähler potential, $F_i = W_i + K_i W / m_{\text{P}}^2$, a subscript i (i^*) denotes derivative with respect to the complex scalar s^i (s^{i^*}), K^{i^*j} is the inverse of K_{ji^*} , and m_{P} is the reduced Planck mass. We will only consider minimal Kähler potential

$$K^{\text{min}} = |S|^2 + |\phi|^2 + |\bar{\phi}|^2 + |H^c|^2 + |\bar{H}^c|^2. \quad (18)$$

The F-term scalar potential then becomes (s is any of the five complex scalars above)

$$V^{\text{min}} = e^{K^{\text{min}}/m_{\text{P}}^2} \left[\sum_s \left| W_s + \frac{W s^*}{m_{\text{P}}^2} \right|^2 - 3 \frac{|W|^2}{m_{\text{P}}^2} \right]. \quad (19)$$

On the trivial path, this scalar potential up to 4th order in $|S|$ is [20]

$$V_{\text{tr}}^{\text{min}} \simeq v_0^4 \left(1 + \frac{1}{2} \frac{|S|^4}{m_{\text{P}}^4} \right) \quad (20)$$

and the effective potential for the standard hybrid case becomes

$$V_{\text{tr}}^{\text{SUGRA}} \simeq V_{\text{tr}}^{\text{min}} + \Delta V \quad (21)$$

with ΔV being the one-loop RC in Eq. (15).

The effective potential on the new smooth path becomes [20]

$$V_{\text{nsm}}^{\text{SUGRA}} \simeq v_0^4 \left(\tilde{V}_{\text{nsm}} + \frac{1}{2} \frac{|S|^4}{m_{\text{P}}^4} \right), \quad (22)$$

where $\tilde{V}_{\text{nsm}} \equiv V_{\text{nsm}}/v_0^4$ with V_{nsm} being the effective potential on the new smooth path for global SUSY, which is constructed [9] numerically.

4.3. Inflationary observables

We can make S real by an appropriate global $U(1)$ R transformation and define the canonically normalized real inflaton field $\sigma \equiv \sqrt{2}\mathcal{S}$. The slow-roll parameters ε , η and the parameter ξ^2 , entering the running of n_s , are given by (see e.g. Ref. [26])

$$\varepsilon \equiv \frac{m_{\text{P}}^2}{2} \left(\frac{V'(\sigma)}{V(\sigma)} \right)^2, \quad \eta \equiv m_{\text{P}}^2 \left(\frac{V''(\sigma)}{V(\sigma)} \right), \quad \xi^2 \equiv m_{\text{P}}^4 \left(\frac{V'(\sigma)V'''(\sigma)}{V^2(\sigma)} \right), \quad (23)$$

where prime denotes derivation with respect to σ and V is either the effective potential $V_{\text{tr}}^{\text{SUGRA}}$ on the trivial path given in Eq. (21), or the effective potential $V_{\text{nsm}}^{\text{SUGRA}}$ on the new smooth path given in Eq. (22).

Numerical simulations show that the system even during the intermediate phase follows basically the new smooth path. The e-foldings from the time when the pivot scale $k_0 = 0.002 \text{ Mpc}^{-1}$ crosses outside the horizon until the end of inflation are [26]

$$N_Q \approx \frac{1}{m_{\text{P}}^2} \int_{\sigma_f}^{\sigma_c} \frac{V_{\text{nsm}}^{\text{SUGRA}}(\sigma)}{V_{\text{nsm}}^{\text{SUGRA}}(\sigma)'} d\sigma + \frac{1}{m_{\text{P}}^2} \int_{\sigma_c}^{\sigma_Q} \frac{V_{\text{tr}}^{\text{SUGRA}}(\sigma)}{V_{\text{tr}}^{\text{SUGRA}}(\sigma)'} d\sigma, \quad (24)$$

where $\sigma_Q \equiv \sqrt{2}\mathcal{S}_Q > 0$ is the value of σ at horizon crossing of the pivot scale k_0 and σ_f the value of σ at the end of the second stage of inflation.

The power spectrum $P_{\mathcal{R}}$ of curvature perturbation at k_0 is given [26] by

$$P_{\mathcal{R}}^{1/2} \simeq \frac{1}{2\pi\sqrt{3}} \frac{V_{\text{tr}}^{\text{SUGRA}}(\sigma_Q)^{3/2}}{m_{\text{P}}^3 V_{\text{tr}}^{\text{SUGRA}}(\sigma_Q)'}. \quad (25)$$

The spectral index n_s , the tensor-to-scalar ratio r , and the running of the spectral index $dn_s/d\ln k$ are [26]

$$n_s \simeq 1 + 2\eta - 6\varepsilon, \quad r \simeq 16\varepsilon, \quad \frac{dn_s}{d\ln k} \simeq 16\varepsilon\eta - 24\varepsilon^2 - 2\xi^2, \quad (26)$$

with ε , η , and ξ^2 evaluated at $\sigma = \sigma_Q$. The number of e-foldings N_Q required for solving the horizon and flatness problems is [27]

$$N_Q \simeq 53.76 + \frac{2}{3} \ln \left(\frac{v_0}{10^{15} \text{ GeV}} \right) + \frac{1}{3} \ln \left(\frac{T_r}{10^9 \text{ GeV}} \right), \quad (27)$$

where T_r is the reheat temperature and should be less than about 10^9 GeV from the gravitino bound [28].

4.4. Monopole Production

Magnetic monopoles are produced at the end of the standard hybrid stage of inflation, where G_{PS} breaks to G_{SM} , via the Kibble mechanism [29]. The initial monopole number density is

$$n_{\text{M}}^{\text{init}} \approx \frac{3p}{4\pi} H_0^3, \quad (28)$$

since H_0^{-1} is the relevant correlation length ($p \sim 1/10$ is a geometric factor). At the end of inflation, the monopole number density becomes

$$n_{\text{M}}^{\text{fin}} \approx \frac{3p}{4\pi} H_0^3 e^{-3\delta N}, \quad (29)$$

where δN is the number of e-foldings of the second inflationary stage.

Dividing $n_{\text{M}}^{\text{fin}}$ by the number density $n_{\text{infl}} \approx V_{\text{tr}}^0/m_{\text{infl}}$ of inflatons produced at the end of inflation (m_{infl} is the inflaton mass), we obtain

$$\frac{n_{\text{M}}}{n_{\text{infl}}} \approx \frac{3p}{4\pi} H_0^3 e^{-3\delta N} \frac{m_{\text{infl}}}{V_{\text{tr}}^0}. \quad (30)$$

This remains fixed until T_{r} , where the relative monopole number density is [30]

$$\frac{n_{\text{M}}}{s} = \frac{n_{\text{M}}}{n_{\text{infl}}} \frac{n_{\text{infl}}}{s} \approx \frac{3p}{16\pi} \frac{H_0 T_{\text{r}}}{m_{\text{P}}^2} e^{-3\delta N} \quad (31)$$

with s being the entropy density.

Taking $n_{\text{M}}/s \lesssim 10^{-30}$ (corresponding [31] to the Parker bound [32]), $T_{\text{r}} \sim 10^9$ GeV, and $H_0 \lesssim 10^{12}$ GeV, we obtain $\delta N \gtrsim 9.2$, which implies that $N_{\text{st}} \lesssim 45$ (N_{st} is the number of e-foldings of the pivot scale k_0 during standard HI). Here, $N_{\text{st}} \ll 45$ and, thus, the present monopole flux is utterly negligible.

4.5. Numerical Results

We put the mass m_A of the color triplet, anti-triplet gauge bosons divided by the gauge coupling constant $g \approx 0.7$ equal to the SUSY GUT scale M_{GUT} and the parameter $p = \sqrt{2}\kappa M/m$ equal to $1/\sqrt{2}$. Also, we take $T_{\text{r}} \simeq 10^9$ GeV saturating the gravitino bound [28] and fix $P_{\mathcal{R}}^{1/2} \simeq 4.85 \times 10^{-5}$ at the scale k_0 from the WMAP normalization [17]. The resulting n_s is plotted [20] against N_{st} and $\alpha = |\langle H^c \rangle|/|\langle \phi \rangle|$ in Fig. 1.

We took $4 \lesssim N_{\text{st}} \lesssim 45$. The lower limit ensures that all the cosmological scales receive perturbations from the first stage of inflation, while the upper limit ensures that the present flux of monopoles in our galaxy does not exceed the Parker bound [32]. Note that, for $\alpha \lesssim 0.2$, λ turns out to be non-perturbative, whereas, for $\alpha \gtrsim 0.7$, the WMAP normalization [17] is not satisfied.

We see that n_s 's below unity are readily obtainable and that the central value $n_s = 0.958$ from WMAP is easily achievable, although spectral indices $n_s \lesssim 0.953$ are not

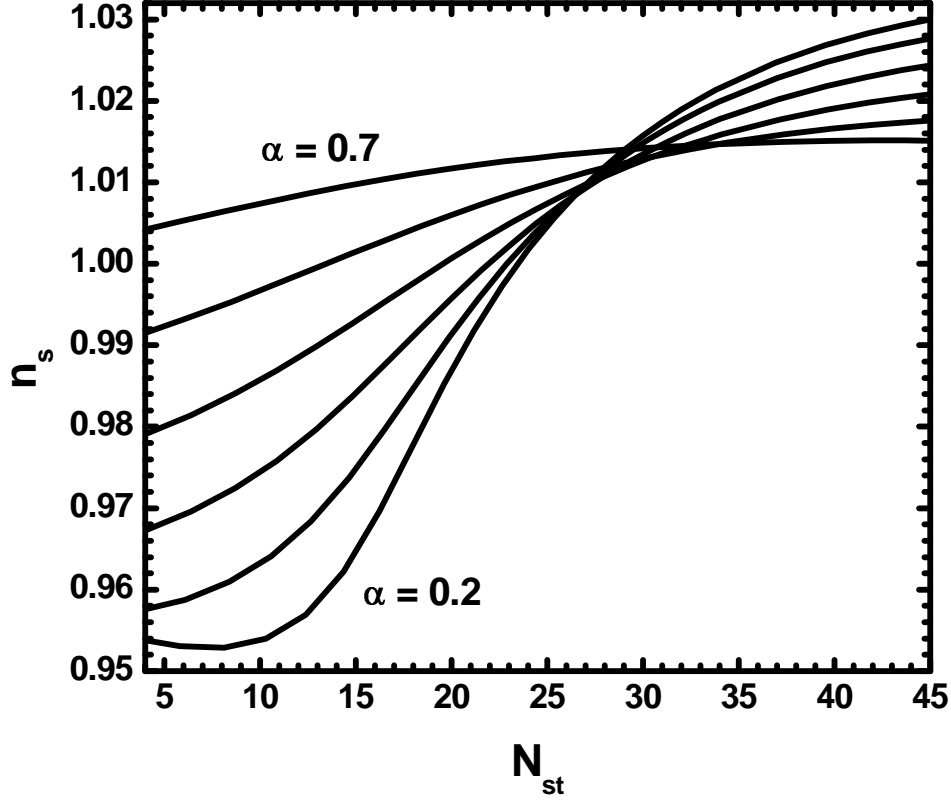


FIGURE 1. Spectral index in standard-smooth HI versus N_{st} in minimal SUGRA for $p = \sqrt{2}\kappa M/m = 1/\sqrt{2}$. The values of the parameter α range from 0.2 to 0.7 with steps of 0.1.

possible. We find that n_s 's in the 95% confidence level range [17]

$$0.926 \lesssim n_s \lesssim 0.99 \quad (32)$$

are obtained only if $N_{\text{st}} \lesssim 21$. So the present monopole flux is expected to be negligible.

The range of the various parameters of the model are [20] $\gamma \simeq (0.17 - 3.43) \times 10^{-3}$, $\kappa \simeq (0.66 - 1.35) \times 10^{-2}$, $\lambda \simeq 0.027 - 0.68$, $M \simeq (2.12 - 2.44) \times 10^{16}$ GeV, $m \simeq (2.8 - 6.6) \times 10^{14}$ GeV, $\sigma_Q \simeq (0.95 - 3.05) \times 10^{17}$ GeV, $\sigma_c \simeq (0.6 - 2) \times 10^{17}$ GeV, $\sigma_f \simeq (4.9 - 9.9) \times 10^{16}$ GeV, $N_Q \simeq 54.1 - 54.5$, $dn_s/d\ln k \simeq -(0.77 - 3.76) \times 10^{-3}$, and $r \simeq (0.7 - 5.3) \times 10^{-5}$.

4.6. Gauge Unification

Cosmology has constrained m to be significantly lower than M_{GUT} , which spoils gauge unification since some fields acquire masses $\sim m$. Actually, the fields with masses $\lesssim M_{\text{GUT}}$ turn out to be too many implying the existence of Landau poles. Also, none of these fields has $\text{SU}(2)_L$ quantum numbers and, thus, the $\text{SU}(2)_L$ gauge coupling constant fails to unify with the other two gauge coupling constants.

Landau poles are avoided by considering the superpotential term $\xi\phi^2\bar{\phi}$ allowed [12] by all the symmetries of the model. This term gives masses $\sim |\xi\langle\phi\rangle|$ to some fields. The second problem is solved by including a single extra superfield $\chi \in (15, 3, 1)$ with mass $m_\chi \approx 8 \times 10^{14}$ GeV and charge $1/2$ under the $U(1)$ R symmetry, which allows it to have just the superpotential term $m_\chi\chi^2/2$. Finally, for gauge unification, we find [20] that $m \gtrsim 4 \times 10^{14}$ GeV, which implies that $\alpha \lesssim 0.5$.

5. CONCLUSIONS

The extension of the SUSY PS model introduced in Ref. [12] in order to solve the m_b problem in SUSY GUTs with YU, such as the simplest SUSY PS model, universal boundary conditions and $\mu > 0$ is a very fruitful framework for constructing HI models. It naturally leads to new shifted and new smooth HI using only renormalizable superpotential terms and avoiding the monopole problem. These variants, however, yield, in minimal SUGRA, too large n_s 's.

This problem can also be resolved within the same model by a two-stage inflationary scenario: the first stage is of the standard and the second of the new smooth hybrid type. Alternatively, we can have semi-shifted HI within the same model. This scenario incorporates cosmic strings contributing to the power spectrum of perturbations. In this case, larger n_s 's are allowed.

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